arXiv:hep-ph/9404220v1 6 Apr 1994

Production and Decay of the Standard Model Higgs Boson at LEP200

Eilam Gross,[†] Bernd A. Kniehl,[‡] and Gustavo Wolf[†]

† Department of Particle Physics

Weizmann Institute of Science, Rehovot 76100, Israel

‡ II. Institut für Theoretische Physik, Universität Hamburg

Luruper Chaussee 149, 22761 Hamburg, Germany

Abstract

We collect and update theoretical predictions for the production rate and decay branching fractions of the Standard Model Higgs boson that will be relevant for the Higgs search at LEP200. We make full use of the present knowledge of radiative corrections. We estimate the systematics arising from theoretical and experimental uncertainties.

1 Introduction

The main mechanism by which the Higgs boson is produced in e^+e^- collisions at LEP is the Higgsstrahlung process where the primary Z boson, *i.e.*, the one that is produced by e^+e^- annihilation, radiates a Higgs boson and then decays into a fermion pair [1, 2].

At LEPI, the primary Z boson is on resonance, $e^+e^- \to Z \to HZ^* \to Hf\bar{f}$ (Bjorken process); see Fig. 1a. As a consequence, the production cross section de-

Figure 1: Feynman diagrams for Higgs-boson production at (a) LEPI and (b) LEP200.

creases dramatically with m_H increasing, and the search potential of LEPI approaches its sensitivity limit, which is probably at $m_H = 65$ GeV or so [3].

At LEP200, operating at centre-of-mass (CM) energies in the range $\sqrt{s} = 170$–200 GeV, the primary Z boson is virtual, while the secondary Z boson, *i.e.*, the one that coexists with the Higgs boson, is close to its mass shell, $e^+e^- \to Z^* \to HZ \to Hf\bar{f}$; see Fig. 1b. As a result, the production cross section is enhanced when $\sqrt{s}$ is large enough to allow both the Higgs and Z bosons to be on their mass shells. LEP200 is, therefore, the suitable machine to look for Higgs bosons, up to about $m_H = 80$–100 GeV, depending on the available CM energy [4].

A dedicated study of the Higgs-boson production cross section at LEP200 does not exist in literature. The Born cross section of $e^+e^- \to Hf\bar{f}$ was calculated analytically by Berends and Kleiss [5]. First-order radiative corrections appropriate to LEPI energies were then calculated using the so-called Improved Born Approximation (IBA) [6] and complemented by a specific heavy-top-quark correction to the ZZH vertex [7]. The full one-loop radiative corrections to the four-point process $Z \to Hf\bar{f}$ were calculated in Ref. [8], where it was shown that, for LEPI energies, the IBA along with the ZZH correction agrees with the fully corrected result for $\Gamma\left(Z \to Hf\bar{f}\right)$ at the level of 1%. However for LEP200 energies, this calculation is no longer adequate because the primary Z-boson is virtual. In principle, one would like to know the full radiative corrections to the five-point process $e^+e^- \to Hf\bar{f}$, but this happens to be very cumbersome and has not been tackled yet. For $\sqrt{s} > m_H + m_Z$, one may resort to the four-point process $e^+e^- \to HZ$, for which the one-loop radiative corrections are known [9, 10]. However, this calculation does not take into account finite-width effects of the secondary Z boson, which is indeed close to its mass shell, but, nevertheless, has an observable width.

In Sect. 2, we make an attempt to predict the five-point process as accurately as possible. To this end, we incorporate in the tree-level result the effects of initial-state electromagnetic bremsstrahlung to second order, including exponentiation of the infrared-sensitive parts. We then plug in the known weak corrections to the four-point process along with experimental information on the Z-boson width. In the limit

of $M_t \gg \sqrt{s}$, the weak corrections are dominated by virtual top-quark contributions. The IBA naturally accounts for most of these terms, except for the one that arises from the one-loop renormalization of the HZZ vertex [7]. The latter must be added by hand. We compare our best estimate with the IBA-type evaluation and find that the two differ appreciably for $\sqrt{s} \gg m_H + m_Z$. Incidentally, the two approaches agree quite well close to threshold. The IBA is applicable also below threshold. On the other hand, the cross section drops rapidly below threshold, so that only $\sqrt{s}$ values down to a few times Γ_Z below $m_H + m_Z$ are relevant phenomenologically. We thus argue that one may still conservatively use the threshold value of the full weak correction to the four-point process in that range below threshold, and estimate the theoretical uncertainty of the result obtained.

In Sect. 3, the Higgs-boson decay branching ratios are re-evaluated and compared with previous results. Many experimental Higgs-search papers quote the Higgs-boson branching ratio to $b\bar{b}$ as 85-87% and that to $\tau^+\tau^-$ as 5-8% [4, 11]. Both branching fractions are of extreme experimental importance. In particular at LEP200, the success of the search for the Higgs boson relies on the b-tagging capability of the experiment. It is obvious that a reliable prediction of $BR\left(H \to b\bar{b}\right)$ is crucial.

The Higgs bosons accessible at LEP200 $(m_H < 100 \text{ GeV})$ are relatively long-lived, with Γ_H of the order of a few MeV, and can thus be taken to be on mass shell in the analysis. The experimentally relevant quantities are thus the total production rate and the various branching fractions including their radiative corrections. Most of the theoretical papers on radiative corrections to Higgs-boson decays consider partial decay widths rather than branching ratios; for a recent review, see Ref. [12]. One purpose of the present paper is to update the Higgs branching ratios relevant for LEP200 making full use of the present knowledge of radiative corrections. The calculations will be shown here in detail.

2 Higgs-Boson Production Cross Section

In this section, we evaluate the radiatively corrected cross section of $e^+e^- \to Hf\bar{f}$ using both the IBA and the full weak corrections to $\sigma(e^+e^- \to HZ)$. Before doing so, we shall briefly review the Born cross section and initial-state bremsstrahlung, which are the core of the analysis.

2.1 Born Approximation

We choose to work in the modified on-mass-shell (MOMS) scheme, in which the Born amplitude is expressed in terms of the Fermi constant, G_F, measured in muon decay so as to suppress large logarithms due to virtual light charged fermions in the radiative corrections. The weak mixing angle is defined as $\sin^2\theta_W = 1 - m_W^2/m_Z^2$. For given m_H and M_t, m_W is fine-tuned such that the perturbative calculation of the

muon lifetime agrees with its high-precision measurement, *i.e.*, such that

$$G_F = \frac{\pi\alpha}{\sqrt{2}\sin^2\theta_W m_W^2} \frac{1}{1 - \Delta r(m_W, m_H, M_t)} \tag{1}$$

is satisfied. Here Δr [13] embodies those radiative corrections to the muon decay width that the Standard Model generates on top of the traditional calculation within the QED-improved Fermi Model. For consistency of our analysis, we use Δr in the one-loop approximation [13]. For a discussion of Δr beyond one loop, we refer to Ref. [14].

In the MOMS scheme, the Born cross section of $e^+e^- \to Hf\bar{f}$, where all fermion flavours $f \neq t$ are summed over, may be written as [5]

$$\sigma_{\mathrm{Born}}^0(s) = \frac{G_F^2(v_e^2 + 1)m_Z^3\Gamma_Z}{96\pi^2 s} \frac{m_Z^2/s}{(1 - m_Z^2/s)^2 + b^2} \int_{x_1}^{x_2} dx\,\mathcal{F}(x), \tag{2}$$

where

$$\mathcal{F}(x) = \frac{(12 + 2a - 12x + x^2)\sqrt{x^2 - a}}{(x - x_p)^2 + b^2}, \tag{3}$$

$v_e = 4\sin^2\theta_W - 1$, $a = 4m_H^2/s$, $b = m_Z\Gamma_Z/s$, $x = 2E_H/\sqrt{s}$, $x_p = 1 + (m_H^2 - m_Z^2)/s$, $x_1 = \sqrt{a}$, $x_2 = 1 + \frac{1}{4}a$, and E_H is the Higgs-boson energy in the laboratory frame. Here the Z-boson propagators are written in the Breit-Wigner form with $m_Z = (91.187 \pm 0.007)$ GeV and $\Gamma_Z = (2.489 \pm 0.007)$ GeV [15]. The integral is performed analytically by means of complex analysis [5].

2.2 Bremsstrahlung

Initial-state corrections to e^+e^- annihilation are of prime importance. These are available to $O(\alpha^2)$ in QED [16, 17]. Here we adopt the formalism of Ref. [17] taking into account real and virtual contributions due to photons and additional e^+e^- pairs. This is achieved by convoluting σ_{Born}^0 with the appropriate radiator function over the full range of center-of-mass energies, $\sqrt{s'}$, accessible after bremsstrahlung. Specifically,

$$\sigma_{\mathrm{Born}}(s) = \int_{x_0}^1 dx\,G(x)\sigma_{\mathrm{Born}}^0(xs), \tag{4}$$

where $x = s'/s$ and $x_0 = m_H^2/s$, taking final-states fermions to be massless. The resummation of infrared-sensitive contributions is accomplished by writing the radiator function, $G(x)$, in an exponentiated form [17],

$$G(x) = \beta(1 - x)^{\beta - 1}\delta^{V+S} + \delta^H(x), \tag{5}$$

where δ^{V+S} and $\delta^H(x)$ are polynomials in $L = \ln(s/m_e^2)$ and $\beta = (2\alpha/\pi)(L - 1)$. Note that here $\alpha = 1/137.035...$ because the emitted photons are real. The term δ^{V+S}

Figure 2: σ_{Born} (dotted line), σ_{IBA} (dashed line), and σ_{full} (solid line) versus $\sqrt{s}$ for $m_H = 80$ GeV assuming $M_t = 165$ GeV.

collects virtual and soft contributions, while δ^H originates from hard bremsstrahlung. For further details, see Ref. [17].

In Fig. 2, σ_{Born} is plotted versus $\sqrt{s}$ for $m_H = 80$ GeV (dotted curve). Numerical results for selected values of $\sqrt{s}$ and m_H are listed in the third column of Table 1.

2.3   Improved Born Approximation

In order to take into account the running of α and virtual heavy-top-quark effects in the evaluation of the Higgs-boson production cross section, it became a common practice among the LEP collaborations to use the IBA [6] complemented by the specific ZZH vertex correction of top origin [7]. In the IBA, the Z-boson and photon propagators get dressed with self-energy insertions, which comprise the leading effects due to light charged fermions and a heavy top quark. These effects may be accommodated by introducing effective parameters, *i.e.*, by substituting $\alpha \to \overline{\alpha} = \alpha(m_Z)$ and $\sin^2 \theta_W \to \sin^2 \overline{\theta}_W = 1 - \cos^2 \overline{\theta}_W = \sin^2 \theta_W + \cos^2 \theta_W \Delta \rho$ in the Born approximation of the on-mass-shell scheme, where

$$1 - \frac{1}{\rho} = \Delta \rho = \frac{3 G_F M_t^2}{8 \pi^2 \sqrt{2}}. \tag{6}$$

Using the relation

$$\rho G_F = \frac{\pi \overline{\alpha}}{\sqrt{2} \sin^2 \overline{\theta}_W \cos^2 \overline{\theta}_W m_Z^2}, \tag{7}$$

which emerges from Eq. (1) by retaining only the dominant terms of Δr, the IBA cross section is obtained from Eq. (2) by substituting $G_F \to \rho G_F \approx (1 + \Delta \rho) G_F$ and $\sin^2 \theta_W \to \sin^2 \overline{\theta}_W$ and including the overall factor $(1 - \frac{8}{3} \Delta \rho)$ [7] to account for the ZZH vertex correction. The result reads

$$\sigma_{\text{IBA}}^0(s) = \frac{G_F^2 (\overline{v}_e^2 + 1) m_Z^3 \Gamma_Z}{96 \pi^2 s} \frac{m_Z^2/s}{(1 - m_Z^2/s)^2 + b^2} (1 + 2 \Delta \rho)(1 - \frac{8}{3} \Delta \rho) \int_{x_1}^{x_2} dx \, \mathcal{F}(x), \tag{8}$$

where $\overline{v}_e = 4 \sin^2 \overline{\theta}_W - 1$. Convolution of σ_{IBA}^0 with G according to Eq. (4) leads to σ_{IBA}, which is potted in Fig. 2, too (dashed curve). Numerical results may be found in the fourth column of Table 1.

2.4   Full One-Loop Radiative Corrections

In the case of $Z \to H f \overline{f}$ relevant for LEPI, the IBA is in reasonable agreement with the full one-loop calculation [8], especially in the upper M_t range. This may be understood by observing that, close to Z-boson peak, the most significant contributions

$\sqrt{s}$ [GeV]	m_H [GeV]	σ_{Born} [pb]	σ_{IBA} [pb]	σ_{full} [pb]
170	60	1.171	1.158	1.139
	70	0.708	0.700	0.692
	80	0.076	0.075	0.075
	90	0.011	0.011	0.011
	100	0.004	0.004	0.004
180	60	1.119	1.106	1.083
	70	0.834	0.825	0.811
	80	0.501	0.496	0.490
	90	0.053	0.053	0.052
	100	0.008	0.008	0.008
190	60	1.016	1.004	0.980
	70	0.821	0.812	0.795
	80	0.610	0.603	0.593
	90	0.365	0.361	0.357
	100	0.038	0.038	0.038
200	60	0.905	0.895	0.871
	70	0.764	0.756	0.738
	80	0.616	0.609	0.597
	90	0.456	0.451	0.443
	100	0.272	0.269	0.265
210	60	0.802	0.793	0.770
	70	0.696	0.688	0.670
	80	0.586	0.579	0.566
	90	0.471	0.465	0.456
	100	0.347	0.343	0.337

Table 1: σ_{Born}, σ_{IBA}, and σ_{full} for several values of $\sqrt{s}$ and m_H GeV assuming $M_t = 165$ GeV.

Figure 3: Distributions of the probability that the $f\bar{f}$ pair produced through $e^+e^- \rightarrow Hf\bar{f}$ at $\sqrt{s} = 180$ GeV has invariant mass $m_{f\bar{f}}$ for $m_H = 70$, 80, and 90 GeV.

arise from loop amplitudes with resonant propagators, the leading terms of which are retained in the IBA. However, this is not necessarily the case at LEP200 energies. In fact, it has been shown [9, 10] that $\sigma(e^+e^- \rightarrow HZ)$ receives sizeable contributions from box diagrams, which do not enter the IBA.

One should keep in mind that in Refs. [9, 10] the Z boson was treated as a stable particle, so that there is no cross section for $\sqrt{s} < m_Z + m_H$. For $\sqrt{s}$ values a few energies above $m_H + m_Z$, this treatment is expected to be a fair one. However, at LEP200, the Higgs detection sensitivity might be pushed to phase space regions where the number of expected Higgs events is very modest, and special care must be exercised. This can be seen in Fig. 3, where the probability that the secondary Z boson is produced with mass $m_{f\bar{f}}$ is shown for $\sqrt{s} = 180$ GeV. This probability is obtained by normalizing the differential cross section,

$$
\frac{d\sigma^0_{\text{Born}}}{dm^2_{f\bar{f}}} \propto \frac{p_{f\bar{f}}\left(3m^2_{f\bar{f}} + p^2_{f\bar{f}}\right)}{\left[(s - m^2_Z)^2 + m^2_Z\Gamma^2_Z\right]\left[\left(m^2_{f\bar{f}} - m^2_Z\right)^2 + m^2_Z\Gamma^2_Z\right]},
\tag{9}
$$

where $p_{f\bar{f}} = \left(\lambda\left(s, m^2_{f\bar{f}}, m^2_H\right)/4s\right)^{1/2}$, with $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2(ab + bc + ca)$, is the $f\bar{f}$ three-momentum in the CM frame. One can clearly see that there is a finite, yet small probability for the Z-boson to be off-shell, especially as m_H approaches the threshold value $\sqrt{s} - m_Z$.

In the following, we suggest a way to incorporate finite-width effects in the one-loop calculation of the Higgs-boson production cross section at LEP200. The calculation of the full radiative corrections to the $2 \rightarrow 3$ process $e^+e^- \rightarrow Hf\bar{f}$ requires the computation of an enormous number of Feynman diagrams and does not exist in literature. However, at least for $\sqrt{s} > m_H + m_Z$, one can take advantage of the fact that the secondary Z-boson is preferably on-shell, so that the radiative corrections factorize approximately into a part connected with the $2 \rightarrow 2$ process $e^+e^- \rightarrow HZ$ and one related to the subsequent Z-boson decay, which are both known. The second part may be included elegantly by using the experimental value of Γ_Z, which is radiatively corrected by nature. Taking into account also initial-state bremsstrahlung, we may write, for $\sqrt{s} > m_H + m_Z$,

$$
\begin{aligned}
\sigma_{\text{full}}(s) &= (1 + 2\mathcal{R}e\Delta_{\text{weak}}(s))\int_{x_0}^{1} dx\, G(x)\sigma^0_{\text{Born}}(xs) \\
&= (1 + 2\mathcal{R}e\Delta_{\text{weak}}(s))\, \sigma_{\text{Born}}(s),
\end{aligned}
\tag{10}
$$

where Δ_{weak} is the finite, gauge-invariant weak correction to $\sigma(e^+e^- \rightarrow HZ)$ as given by Eq. (4.8) of Ref. [9]. One might wonder whether the weak correction term should

Figure 4: $\sigma_{\text{full}}/\sigma_{\text{IBA}}$ versus $\sqrt{s}$ for $m_H = 80$ GeV assuming $M_t = 165$ GeV.

be evaluated at xs and included as part of the integrand. We have chosen not to do so. However, this question becomes irrelevant by noticing that this modification changes the cross section by less than 0.1%.

For $\sqrt{s} < m_H + m_Z$, one faces the problem that Δ_{weak} is not defined, since the secondary Z boson is pushed from its mass shell. For the sake of continuity, we propose to use Eq. (10) with Δ_{weak} evaluated at $(m_H + m_Z)^2$ instead of s. This procedure turns out to be a conservative one as will be shown below.

In Fig. 2, σ_{full} (solid line) is compared with σ_{Born} and σ_{IBA}. Numerical values are listed in the last column of Table 1. We observe that the IBA overshoots the full one-loop calculation for $\sqrt{s} \gg m_H + m_Z$. To elaborate this point, we show in Fig. 4 the ratio $\sigma_{\text{full}}/\sigma_{\text{IBA}}$ as a function of $\sqrt{s}$. It is clearly seen that the IBA rapidly deteriorates as $\sqrt{s}$ increases. On the other hand, the ratio gets close to one as $\sqrt{s}$ approaches $m_H + m_Z$, the value at threshold being 0.995. For $\sqrt{s} < m_H + m_Z$, the ratio is independent of $\sqrt{s}$, since we continue to use the threshold value of Δ_{weak} in the evaluation of σ_{full}. The corrections implemented in the IBA are independent of $\sqrt{s}$ anyway. Thus, σ_{IBA} is perfectly well defined theoretically also below threshold. Nevertheless, for the sake of a continuous description, we are in favour of using σ_{full}, provided that $\sqrt{s}$ is not more than a few times Γ_Z below $m_H + m_Z$. This attitude is also conservative from the experimental point of view, since σ_{IBA} might overestimate the true production rate. However, since both approaches agree to the level of 0.5%, we do not anticipate a great theoretical uncertainty. However, a firm conclusion concerning the virtue of this approximation can be drawn only from a full one-loop calculation of $\sigma\left(e^+e^- \to Hf\bar{f}\right)$. In Fig. 2, we assumed $m_H = 80$ GeV. In Fig. 5, we show the $\sqrt{s}$ dependence of σ_{full} also for other values of m_H. In practice, $\sqrt{s}$ will be fixed at some value between 170 and 200 GeV. It is then useful to know the m_H dependence of σ_{full}, which is shown in Fig. 6.

2.5 Systematics

The accuracy of the predicted value of σ_{full} is primarily limited by the errors on the input parameters. Apart from m_H, these are m_Z, Γ_Z, and M_t. m_Z and Γ_Z have been measured at LEP to high accuracy, $m_Z = (91.187 \pm 0.007)$ GeV and $\Gamma_Z = (2.489 \pm 0.007)$ GeV [15]. Recent global analyses [15] of LEP data suggest $M_t = (166^{+17+19}_{-19-22})$ GeV via loop effects. Recently, the D0 Collaboration at the Fermilab Tevatron has announced a lower limit of 131 GeV on M_t [18]. Therefore, $M_t = (165 \pm 35)$ GeV covers the most probable M_t range.

Figure 5: σ_{full} versus $\sqrt{s}$ for $m_H = 70, 80, 90,$ and 100 GeV assuming $M_t = 165$ GeV.

Figure 6: σ_{full} versus m_H for $\sqrt{s} = 170$, 180, 190, and 200 GeV assuming $M_t = 165$ GeV.

Changing m_Z (Γ_Z) by $\pm 2\sigma$ shifts σ_{full} by at most 0.6% (0.5%); the maximum effect occurs when $\sqrt{s} = m_H + m_Z$ ($\sqrt{s} < m_H + m_Z$). The variation of σ_{full} with M_t is studied in Fig. 7 for typical LEP200 conditions, $\sqrt{s} = 180$ GeV and $m_H = 70$ GeV. We see that the present uncertainty in M_t induces a systematic error of $\pm 0.5\%$ in σ_{full}. Since the one-loop corrections to the Higgs-boson production rate at LEP200 are relatively modest in the MOMS scheme, below 4% in magnitude, we expect that the theoretical uncertainty due to unknown higher orders is insignificant. In summary, we estimate the total systematic error on the Higgs-boson production cross section to be of the order of 0.9%.

3 Higgs-Boson Decay Branching Ratios

In this section, we study in detail the branching ratios of the Higgs boson that are relevant at LEPI and LEP200. We start by reviewing the tree-level results. We then analyze the influence of radiative corrections. Finally, we estimate the systematic errors involved in the calculations.

3.1 Born Approximation

The Higgs boson couples directly to fermions thereby generating their masses. At tree level, the coupling strength is $2^{1/4} G_F^{1/2} m_f$ and the $H \to f\bar{f}$ decay width is

$$\Gamma_0 \left(H \to f\bar{f} \right) = \frac{N_c G_F m_H m_f^2 \beta_f^3}{4\pi\sqrt{2}}, \tag{11}$$

where $N_c = 1$ (3) for leptons (quarks) and $\beta_f = \sqrt{1 - 4m_f^2/m_H^2}$ is the velocity of f in the CM frame. Thus, the tree-level branching fractions are

$$BR \left(H \to f\bar{f} \right) = \frac{N_c m_f^2 \beta_f^3}{\sum_{f' \neq t} N_c m_{f'}^2 \beta_{f'}^3}. \tag{12}$$

Using $m_b = 4.7$ GeV, $m_c = 1.45$ GeV [19], and $m_\tau = 1.777$ GeV, one finds $BR \left(H \to b\bar{b} \right) = 87\%$, $BR \left(H \to c\bar{c} \right) = 9\%$, and $BR \left(H \to \tau^+\tau^- \right) = 4\%$. These lowest-order estimates are subject to electroweak and QCD corrections. Moreover, higher-order decay channels need to be taken into account.

Figure 7: σ_{full} versus M_t for $\sqrt{s} = 180$ GeV and $m_H = 70$ GeV.

9

3.2 Higgs-Boson Decays to Two Electroweak Bosons

We start by considering the decay of the Higgs boson to four fermions via a pair of virtual W bosons. This channel becomes relevant for LEP200 as soon as $m_H > m_W$, so that one W boson can get on its mass shell. Its partial decay width can be written as [20]

$$
\Gamma(H \to W^*W^*) = \int_{\sqrt{s_+}+\sqrt{s_-}\leq m_H} \frac{ds_+ds_-}{\pi^2} \frac{s_+\Gamma_W/m_W}{\left(s_+ - m_W^2\right)^2 + m_W^2\Gamma_W^2}
$$
$$
\times \frac{s_-\Gamma_W/m_W}{\left(s_- - m_W^2\right)^2 + m_W^2\Gamma_W^2}\Gamma\left(m_H^2, s_+, s_-\right), \tag{13}
$$

where

$$
\Gamma\left(m_H^2, s_+, s_-\right) = \frac{3G_F m_W^4}{2\pi\sqrt{2}m_H^3}\sqrt{\lambda\left(m_H^2, s_+ s_-\right)}\left(1 + \frac{\lambda\left(m_H^2, s_+ s_-\right)}{12s_+ s_-}\right), \tag{14}
$$

where $\lambda(a, b, c)$ is defined below Eq. (9) and we use the experimental values $\Gamma_W = 2.12$ GeV and $m_W = 80.22$ GeV [21]. The formula for the $H \to Z^*Z^*$ decay width emerges from Eq. (13) by substituting m_W and Γ_W by m_Z and Γ_Z, respectively, and including the factor $1/2$ to account for identical-particle symmetrization.

The $H \to \gamma\gamma$ [2, 22] and $H \to \gamma Z$ [23] decays proceed through W-boson and charged-fermion loops and are generally less significant for the Higgs search at LEP 200. QCD corrections to their partial widths are well under control [24, 25].

3.3 QCD and Electroweak Corrections

In the case of $H \to q\bar{q}$, it is important to include QCD corrections to Eq. (11) [26, 27]. In fact, when the pole mass, M_q, is used as a basic parameter, these corrections contain large logarithms of the form $(\alpha_s/\pi)^n \ln^m\left(m_H^2/M_q^2\right)$, with $n \geq m$. Appealing to the renormalization-group equation, these logarithms may be absorbed completely into the running quark mass, $m_q(\mu)$, evaluated at scale $\mu = m_H$. In this way, these logarithms are resummed to all orders and the perturbation expansion converges more rapidly. This observation gives support to the notion that the $Hq\bar{q}$ Yukawa couplings are controlled by the running quark masses.

The values of M_q may be estimated from QCD sum rules. In our analysis, we use $M_c = (1.45 \pm 0.05)$ GeV and $M_b = (4.7 \pm 0.2)$ GeV [19]. To obtain $m_q(m_H)$, we proceed in two steps. Firstly, we evaluate $m_q(M_q)$ from

$$
m_q(M_q) = \frac{M_q}{1 + (4/3)\alpha_S(M_q)/\pi + K(\alpha_S(M_q)/\pi)^2}, \tag{15}
$$

with [28]

$$
K \approx 16.11 - 1.04 \sum_{i=1}^{n_F-1}\left(1 - \frac{M_i}{M}\right), \tag{16}
$$

Figure 8: μ dependence of (a) $m_c(\mu)$ and (b) $m_b(\mu)$ evaluated to $O\left(\alpha_S^2\right)$ assuming $M_c = 1.45$ GeV, $M_b = 4.7$ GeV, and $\alpha_S(m_Z) = 0.123$.

where the sum extents over all quark flavours with $M_i < M_q$. Specifically, $K_c = 13.2$ and $K_b = 12.3$. Secondly, we determine $m_q(m_H)$ via the scaling law

$$m_q(m_H) = m_q(M_q)\frac{c_q(\alpha_S(m_H)/\pi)}{c_q(\alpha_S(M_q)/\pi)}, \tag{17}$$

where

$$c_c(x) = \left(\frac{25}{6}x\right)^{12/25}(1 + 1.014x + 1.389x^2), \tag{18}$$

$$c_b(x) = \left(\frac{23}{6}x\right)^{12/23}(1 + 1.175x + 1.501x^2). \tag{19}$$

Throughout our analysis, we evaluate α_S from the two-loop formula [21],

$$\alpha_S(\mu) = \frac{12\pi}{(33 - 2n_F)\ln\left(\mu^2/\Lambda_{(n_F)}^2\right)}\left[1 - \frac{6(153 - 19n_F)}{(33 - 2n_F)^2}\frac{\ln\left[\ln\left(\mu^2/\Lambda_{(n_F)}^2\right)\right]}{\ln\left(\mu^2/\Lambda_{(n_F)}^2\right)}\right], \tag{20}$$

where n_F is the number of quark flavours active at scale μ and $\Lambda_{(n_F)}$ is the appropriate asymptotic scale parameter. We fix $\Lambda_{(5)}$ by requiring that $\alpha_S(m_Z) = 0.123$ [29] and determine $\Lambda_{(4)}$ from the condition that $\alpha_s(\mu)$ be continuous at the flavour threshold $\mu = M_b$. Using the above value of M_b, we find $\Lambda_{(4)} = 0.416$ GeV and $\Lambda_{(5)} = 0.296$ GeV. The scale dependences of m_c and m_b to $O\left(\alpha_S^2\right)$ are illustrated in Fig. 8 and Table 2.

m_H [GeV]	$m_b(m_H)$ [GeV]	$m_c(m_H)$ [GeV]
50	2.89	0.51
60	2.84	0.50
70	2.80	0.50
80	2.77	0.49
90	2.74	0.48
100	2.71	0.48

Table 2: $m_b(m_H)$ and $m_c(m_H)$ evaluated to $O\left(\alpha_S^2\right)$ for several values of m_H assuming $M_c = 1.45$ GeV, $M_b = 4.7$ GeV, and $\alpha_S(m_Z) = 0.123$.

11

The QCD corrections to $\Gamma\left(H \to q\bar{q}\right)$ are known up to $O\left(\alpha_S^2\right)$ for $q \neq t$. In the $\overline{\text{MS}}$ scheme, the result is [12]

$$
\begin{aligned}
\Gamma\left(H \to q\bar{q}\right) &= \frac{3 G_F m_H m_q^2}{4\pi\sqrt{2}} \left[\left(1 - 4\frac{m_q^2}{m_H^2}\right)^{3/2} + \frac{\alpha_S}{\pi}\left(\frac{17}{3} - 40\frac{m_q^2}{m_H^2} + O\left(\frac{m_q^4}{m_H^4}\right)\right)\right. \\
&\quad \left. + \left(\frac{\alpha_S}{\pi}\right)^2\left(K_2 + O\left(\frac{m_q^2}{m_H^2}\right)\right) + O\left(\left(\frac{\alpha_S}{\pi}\right)^3\right)\right],
\end{aligned}
\tag{21}
$$

where $K_2 \approx 35.9399 - 1.3586 n_F$ [27] and it is understood that α_S and m_q are to be evaluated at $\mu = m_H$. We note in passing that Eq. (21) may be translated into the on-mass-shell scheme by using the above relation between M_q and $m_q(m_H)$ [30]. However, appealing to the general notion that the resummation of large logarithms is automatically implemented by the $\overline{\text{MS}}$ evaluation using the appropriate scale, we express a preference for the use of Eq. (21). The difference between these two evaluations is extremely small [30], which indicates that the residual uncertainty due to the lack of knowledge of the $O\left(\alpha_S^2 m_q^2/m_H^2\right)$ and $O\left(\alpha_S^3\right)$ terms is likely to be inconsequential for practical purposes.

The hadronic width of the Higgs boson receives contributions also from the $H \to gg$ channel, which is mediated by massive-quark triangles, and related higher-order processes. The respective partial width is well approximated by [31]

$$
\Gamma\left(H \to gg(g), gq\bar{q}\right) = \Gamma\left(H \to gg\right)\left(1 + \frac{\alpha_S(m_H)}{\pi}\left(\frac{95}{4} - \frac{7}{6}n_F\right)\right),
\tag{22}
$$

where, in the m_H range of current interest, $n_F = 5$ and [32]

$$
\Gamma(H \to gg) = \frac{\alpha_S^2(m_H) G_F m_H^3}{36\pi^2\sqrt{2}}\left(1 + \frac{7}{60}\frac{m_H^2}{M_t^2} + O\left(\frac{m_H^4}{M_t^4}\right)\right).
\tag{23}
$$

Finally, we discuss the one-loop electroweak corrections to the $H \to f\bar{f}$ decay rates [33, 34]. In the m_H range under consideration, these may be incorporated by multiplying Eq. (11) with [33]

$$
\begin{aligned}
\mathcal{K} &= \left\{1 + \frac{\alpha}{\pi}\frac{3}{2}Q_f^2\left(\frac{3}{2} - \ln\frac{m_H^2}{m_f^2}\right) + \frac{G_F}{8\pi^2\sqrt{2}}\left[C_f M_t^2 + m_W^2\left(3\frac{\ln\cos^2\theta_W}{\sin^2\theta_W} - 5\right)\right.\right. \\
&\quad \left.\left. + m_Z^2\left(\frac{1}{2} - 3\left(1 - 4\sin^2\theta_W|Q_f|\right)^2\right)\right]\right\},
\end{aligned}
\tag{24}
$$

where Q_f is the electric charge of f (in units of the positron charge), $C_b = 1$, and $C_f = 7$ for $f \neq t, b$. Note that for $f = b$ the M_t dependence is strongly reduced due to a cancellation between universal self-energy diagrams and specific triangles involving virtual top quarks. In the case of $f \neq t, b$, also the two-loop corrections of

12

Figure 9: Branching fractions of the Higgs boson in the m_H window relevant for LEPI and LEP200. All radiative corrections discussed in the text are included.

$O\left(\alpha_S G_F M_t^2\right)$ are known [35]. They screen somewhat the one-loop M_t dependence, so that effectively $C_f = 7 - 2(\pi/3 + 3/\pi)\alpha_S \approx 7 - 4\alpha_S$. Electroweak corrections to the fermionic decay rates of the Higgs boson are relatively modest in the m_H range accessible at LEP200, the maximum effect being 1.2% in the $\tau^+\tau^-$ channel, 0.3% for $c\bar{c}$, and 0.6% for $b\bar{b}$.

We are now in a position to compute accurately all the Higgs-boson branching ratios that are of phenomenological relevance at LEPI and LEP200. Our final results are summarized in Fig. 9 and Table 3. To appreciate the effect of the radiative corrections to the decay widths on the branching fractions, we contrast in Table 4 our final results for $m_H = 80$ GeV with the evaluations based on the tree-level decay rates along with the quark pole masses according to Eq. (12). The sole effect of running the quark masses in the tree level formulae is also shown. One can see that the latter makes up the dominant effect on the resulting branching ratios.

3.4 Systematics

We do not include the $H \rightarrow s\bar{s}$ channel in our analysis, since its partial width is greatly suppressed by the smallness of $m_s(m_H)$. However, uncertainties in M_c, M_b, M_t, and $\alpha_S(m_Z)$ are found to be significant. Their maximum effects on the branching ratios are investigated in Table 5 for $m_H = 80$ GeV. The numbers are very similar for other values of m_H. The total systematics is obtained by combining the individual errors in quadrature. The present situation is visualized in Fig. 10.

4 Conclusion

In this paper, we have collected and updated the theoretical predictions for the Higgs-boson production rate and decay branching fractions appropriate to LEP200 conditions. Our analysis of $\sigma\left(e^+e^- \rightarrow Hf\bar{f}\right)$ includes initial-state bremsstrahlung to second order with exponentiation, finite-width effects, and the full one-loop weak corrections to the underlying $2 \rightarrow 2$ process, $e^+e^- \rightarrow HZ$. We showed that the popular Improved Born Approximation supplemented with the same initial-state corrections deviates appreciably from the evaluation with the full weak corrections, especially at energies far above the threshold of on-shell HZ production. We also suggested a way to implement weak corrections below this threshold. We have assigned a 0.9%

Figure 10: Branching ratios of the Higgs-boson decays to (a) $b\bar{b}$, (b) $\tau^+\tau^-$, $c\bar{c}$, W^*W^*, (c) gg, $\gamma\gamma$, γZ, and Z^*Z^* versus m_H. The bands indicate the systematic errors.

13

m_H [GeV]	bb	$\tau^+\tau^-$	$c\bar{c}$	gg	W^*W^*	Z^*Z^*	$\gamma\gamma$
50	88.3	8.7	2.8	0.2	-	-	-
60	87.9	9.0	2.8	0.3	-	-	-
70	87.6	9.2	2.8	0.4	-	-	-
80	87.1	9.5	2.8	0.4	0.1	-	0.1
90	86.7	9.6	2.8	0.6	0.2	-	0.1
100	85.3	9.7	2.7	0.7	1.3	0.1	0.2
110	80.8	9.4	2.6	0.8	5.7	0.5	0.2

Table 3: Branching fractions (in %) of the Higgs boson in the m_H window relevant for LEPI and LEP200. All radiative corrections discussed in the text are included.

Decay Mode	Tree-Level	$m_q \to m_q(m_H)$	Full
$H \to bb$	87.3	85.5	87.1
$H \to \tau^+\tau^-$	4.2	11.8	9.5
$H \to c\bar{c}$	8.5	2.7	2.8
$H \to gg$	-	-	0.4
$H \to W^*W^*$	-	-	0.1
$H \to \gamma\gamma$	-	-	0.1

Table 4: Branching fractions (in %) of an 80 GeV Higgs boson evaluated from Eq. (12) with M_q and $m_q(m_H)$, respectively, and evaluated including all channels and radiative corrections discussed in the text.

Decay Mode	BR [%]	$M_b = (4.7 \pm 0.2)$ GeV	$M_c = (1.45 \pm 0.05)$ GeV	$\alpha_S(m_Z) = 0.123 \pm 0.006$	$M_t = (165 \pm 35)$ GeV
$H \to b\bar{b}$	87.1	$+1.2$ -1.3	$+0.4$ -0.3	$+0.0$ -0.1	$+0.1$ -0.0
$H \to c\bar{c}$	2.8	$+0.3$ -0.3	$+0.4$ -0.4	$+0.9$ -1.1	$+0.0$ -0.0
$H \to gg$	0.4	$+0.1$ -0.0	$+0.1$ -0.1	$+0.1$ -0.0	$+0.0$ -0.0
$H \to \tau^+\tau^-$	9.4	$+1.0$ -0.8	$+0.1$ -0.0	$+1.0$ -0.8	$+0.1$ -0.0

Table 5: Effects of the uncertainties in M_q and $\alpha_S(m_Z)$ on the various Higgs branching ratios for $m_H = 80.0 \ GeV$.

systematical error to our results arising mainly from the uncertainties in m_Z, Γ_Z, and M_t.

In our analysis of the Higgs-boson branching fractions, we took into account two-loop QCD corrections to the hadronic widths, one-loop electroweak corrections to the fermionic widths, as well as the contributions from the $\gamma\gamma$, γZ, Z^*Z^*, and W^*W^* channels. The branching ratios of $H \to b\bar{b}$ and $H \to \tau^+\tau^-$ were found to differ slightly from the commonly quoted numbers. *E.g.*, for $m_H = 80$ GeV, we obtained $BR\left(H \to b\bar{b}\right) = (87.1\pm1.3)\%$ and $BR(H \to \tau^+\tau^-) = (9.4^{+1.4}_{-1.1})\%$. We also estimated the systematics on the branching fractions, which should be relevant already for the determination of the m_H lower bound at LEPI.

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